Is the $\xi(2.2)$ a Higgs Boson?

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The question whether the $\xi(2.2)$ recently discovered in $J/\psi \to \gamma + \xi$ can be a Higgs boson in a nonminimal Higgs scheme is discussed. It is pointed out that a search for ξ in B decays provides a critical test.

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With the experimental observation of the intermediate vector bosons (W, Z), the $SU(2) \otimes U(1)$ non-Abelian gauge theory of the weak interactions has received striking confirmation. However, one crucial aspect of that theory, the Higgs mechanism for generating the masses of the W and Z, has not received similar experimental confirmation; no Higgs boson has been found. The Mark III detector group has reported[1] the observation of a particle $\xi(2.2)$ of mass 2.2 GeV and width less than about 30 MeV (resolution) in decays $J/\psi \to \gamma + \xi$, with dominant two-body decay modes $\xi \to K\bar{K}$. The decay mode $K_S K_S$ is seen, so that the spin is even. The width is much less than expected for an ordinary 2-GeV hadron. The η_c, which is narrow because of the Okubo-Zweig-Iizuka suppression of $c\bar{c}$ decays, has already been identified at 2.98 GeV. It is also probably too narrow for a 2-GeV glueball; and a glueball, being a flavor singlet, would be expected to decay into π's at least as often as into K's. The decay $V \to \gamma + H$, where V is a $Q\bar{Q}$ bound state with Q a heavy quark, was suggested[2] some time ago as a favorable place to search for a Higgs boson of mass $m_H < m_V$. And $K\bar{K}$ is the expected dominant two-body decay mode for a Higgs boson with mass in the 2-GeV range. Thus if the spin can be established to be zero and the width remains consistent with resolution, the $\xi(2.2)$ should be considered[3] as a candidate for a Higgs boson.

There are difficulties with the identification of the $\xi(2.2)$ with the physical neutral Higgs boson of the minimal (one Higgs doublet) standard model.

(1) There is the Linde-Weinberg argument,[4] based on the stability of the broken-symmetry vacuum state, that m_H should be greater than 7 to 10 GeV.

(2) The product of the branching ratios for $J/\psi \to \gamma + \xi$ and $\xi \to K^+K^-$ is measured[1] as

$$B(\psi \to \gamma + \xi) B(\xi \to K^+K^-) = (8.0 \pm 2.6) \times 10^{-5}. \quad (1)$$

The branching ratio for $\xi \to K^+K^-$ must be less than $\frac{1}{2}$ ($\xi \to K^0\bar{K}^0$ is equally likely; $\xi \to K_S K_S$ is observed).[1] If the observed decays are $H \to s\bar{s} \to K\bar{K}$, the 2.2-GeV $s\bar{s}$ pair will hadronize into many channels, so that $B(\xi \to K^+K^-)$ will be substantially less than $\frac{1}{2}$. The branching ratio for $J/\psi \to \gamma + H$ can be calculated[2] in the minimal standard model:

$$B_0(V \to \gamma + H) = \frac{G m_Q^2}{\sqrt{2}\pi\alpha}\left(1 - \frac{m_H^2}{m_V^2}\right) B(V \to e^- e^+), \quad (2)$$

for V a nonrelativistic $Q\bar{Q}$ bound state. With use of the experimental branching ratio for $J/\psi \to e^- e^+$ ($\simeq 0.074$) and $m_c \simeq 1.5$ GeV, this gives $B_0(\psi \to \gamma + H) \simeq 2.9 \times 10^{-5}$. Thus the calculated upper limit for (1) is at least an order of magnitude smaller than the experimental value. Put another way, the experimental result determines an enhancement factor $[B(\xi \to K^+K^-) = b]$ as follows:

$$\frac{B(\psi \to \gamma + \xi)}{B_0(\psi \to \gamma + H)} = \frac{(8.0 \pm 2.6) \times 10^{-5}}{2.9 \times 10^{-5}} \frac{1}{b} \equiv \frac{r}{b}, \quad (3)$$

with $r_n = r$(nominal) $= 2.8$, or $r_c = r$(90% C.L.) $= 1.3$ (C.L. denotes confidence level).

(3) The decay $\xi \to \mu^- \mu^+$ has not been observed.[1] The simple expectation is

$$\frac{\Gamma(H \to \mu\bar{\mu})}{\Gamma(H \to K\bar{K})} > \frac{\Gamma(H \to \mu\bar{\mu})}{\Gamma(H \to s\bar{s})} \simeq \frac{1}{3} \frac{m_\mu^2}{m_s^2}, \quad (4)$$

where m_s is a Lagrangian, or "current quark," mass, whose value is not well determined; but one expects this lower bound to be in the range of a few to twenty percent. If the branching ratio b is substantially less than 1, then the signal for $\xi \to \mu\bar{\mu}$ should not be much less than that for $\xi \to K\bar{K}$.

These difficulties seem to rule out the identification of the $\xi(2.2)$ with the single physical neutral Higgs boson of the minimal standard model. I therefore consider whether these problems can be resolved by enlarging the Higgs sector. I will restrict the generalization to the inclusion of additional $SU(2)_w$ (complex) doublets of Higgs bosons. This is a sufficient[5] condition to maintain the empirically successful result of the minimal model, $m_W^2/m_Z^2\cos^2\theta_w \simeq 1$. Furthermore, the experimental absence of flavor-changing neutral currents leads to the requirement[6] that all fermions of a given charge couple to only one Higgs doublet. Higgs sectors consisting of more than one doublet have a variety of theoretical motivations, predating the observation of the $\xi(2.2)$, and various aspects of some such models have been discussed in the literature. It appears that when one goes beyond the minimal one-Higgs-doublet model, no lower bound can be placed on the Higgs boson masses, except possibly one of the physical neutral scalars.[4,7] In the minimal model, the couplings of the Higgs bosons to the other particles are determined. A number of papers have discussed the possibility of enhancing some of these couplings in a nonminimal model. It was already noted[2,7-9] that discovery of a Higgs boson with a branching ratio for $V \to \gamma + H$ substantially greater than the theoretical value (2) for the minimal model would be a signal for an extended Higgs sector.

We consider N_H doublets Φ_I. The neutral member of Φ_I has vacuum expectation value (VEV) $v_I/\sqrt{2}$, with $\sum v_I^2 = v^2$ and $m_W^2 = (g^2/4)v^2$. To minimize the Higgs potential and diagonalize the quadratic part is in general quite complicated, but in the important special case that one v_{I_0} is much less than any other, the corresponding Φ_{I_0} does not mix with the other Φ_I in the quadratic part of $V(\Phi)$, and the mass of its neutral scalar component is proportional to v_{I_0}. So we have an arbitrary small parameter at our disposal,

$$\beta = v_{I_0}/v \ll 1 \text{ (arbitrary)}, \tag{5}$$

and a neutral scalar Higgs, call it h^0, which is typically much lighter than the other Higgs particles (H_I), $m_h \sim \beta m_H$.

The Yukawa coupling constant of the Ath-generation fermion of charge Q $(-1, +\frac{2}{3}, -\frac{1}{3})$, coupled to Φ_I in the fermion-Higgs Lagrangian, to the neutral scalar H_J^0, is

$$\frac{g}{2}\frac{m_A}{m_W}\left(\frac{v}{v_I}\right)\Omega(H^0)_{IJ}, \tag{6}$$

where $\Omega(H^0)$ is the orthogonal matrix which arises in diagonalizing the neutral scalar sector of the quadratic part of $V(\Phi)$. Since $\Omega_{I_0J} = \delta_{I_0J} + O(\beta)$, the fermions of charge Q, which are coupled to Φ_{I_0}, have enhanced coupling to the light h^0, by the factor $(v/v_0)\Omega(H^0)_{I_0I_0} = 1/\beta$. The couplings of fermions of other charges to h^0 are suppressed by the factor $(v/v_I)\Omega(H^0)_{II_0} = O(\beta)$, since $I \neq I_0$.

Note that the enhancement factor is precisely $1/\beta$ (for $\beta \ll 1$), while the suppression factors are only $O(\beta)$; the numerical factors depend on the arbitrary parameters of the potential.

Returning to the specific case of the $\xi(2.2)$, we identify it with the light h^0. This resolves problem (1). We couple the $Q=\frac{2}{3}$ quarks to Φ_{I_0}. Then $1/\beta^2$ is the required enhancement factor (3),

$$1/\beta^2 = B(\psi \to \gamma + \xi)/B_0(\psi \to \gamma + h^0) = r/b \gtrsim 10, \tag{7}$$

and probelm (2) is resolved. If we also couple the $Q=-\frac{1}{3}$ quarks to Φ_{I_0}, and the leptons to a different Φ_I, then problem (3) is resolved. Note that two Higgs doublets are sufficient for this scheme. But this scheme implies that the branching ratio for $\Upsilon \to \gamma + h^0$ is enhanced by the same factor $(1/\beta^2)$ as $\psi/J \to \gamma + h^0$. The branching ratio computed from (2), with $m_b \simeq 4.9$ GeV and $B_{ee} \simeq 0.03$, is

$$B(\Upsilon \to \gamma + h^0)_{\text{th}} \simeq 2.4\times10^{-4} \times 1/\beta^{-2} \gtrsim 2\times10^{-3}. \tag{8}$$

There is an experimental limit from the CUSB detector group.[10] In this mass range

$$B(\Upsilon \to \gamma + h^0)_{\text{expt}} < 10^{-3} \text{ (90\% C.L.)}. \tag{9}$$

There are enough experimental uncertainties, r and b in (7) and B_{ee} in (8), that within a couple of standard deviations (8) and (9) might coexist, leaving a simple two-doublet model viable. However, we will continue under the assumption that (8) and (9) require that the $Q=-\frac{1}{3}$ quarks not be coupled to Φ_{I_0}. In this case the couplings of h^0 to $Q=-\frac{1}{3}$ quarks and to leptons are both suppressed, and problem (3) is open again. Another problem is the following. One can compute[2] the ratio

$$\frac{\Gamma(h^0 \to gg)}{\Gamma(h^0 \to s\bar{s})} \simeq \frac{4\alpha_s^2}{27\pi^2}\left(\frac{m_H^2}{m_s^2}\right)\frac{1}{\beta^2}\,\frac{1}{O(\beta^2)}, \tag{10}$$

where the first $1/\beta^2$ comes from the enhancement of the h^0 coupling to the virtual $Q=\frac{2}{3}$ quark loop in $h^0 \to gg$ and the $O(\beta^2)$ comes from the suppression of the $h^0 \to s\bar{s}$ coupling. If we take $O(\beta^2) = \beta^2 \lesssim \frac{1}{10}$, then this branching ratio is $\gtrsim 1.5$. Since two gluons do not hadronize preferentially into $K\bar{K}$, this is inconsistent with $K\bar{K}$ being the dominant two-body decay mode. These two problems rule out models in which the suppression factors for $h^0 \to s\bar{s}$ and $h^0 \to \mu\bar{\mu}$ are the same and both $\lesssim \beta$. This is in fact the situation in a two-doublet model if the $Q=-\frac{1}{3}$ quarks are not coupled to Φ_{I_0}. But, as we noted above, the suppression factors are strongly model dependent, and one might guess that by further extending the Higgs sector to $N_H \geqslant 3$, we might be able to evade these problems. Indeed, it is not hard to see how to accomplish this in an $N_H = 3$ model. Couple the leptons to Φ_1, the $Q = -\frac{1}{3}$ quarks to Φ_2, and the $Q = \frac{2}{3}$ quarks to Φ_3, and take $v_3/v \sim \epsilon$, $v_2/v \to \sqrt{\epsilon}$, $v_1/v \sim 1$. Then the fermion-h^0 couplings are (for $\epsilon \ll 1$)

$$\sim (1/\epsilon) h^0 \bar{u} u, \quad \sim \epsilon (1/\sqrt{\epsilon}) h^0 \bar{d} d, \quad \sim \epsilon \epsilon^{1/2} h^0 \bar{l} l. \tag{11}$$

Again, the coefficient of $1/\epsilon$ is precisely 1; the coefficients of the other terms depend on the arbitrary parameters of V.

There is another class of processes to consider, which has the virtue that an adverse consequence cannot be evaded by increasing N_H or adjusting the parameters in V. That is the class of neutral flavor-changing processes that go by way of loops of $Q = \frac{2}{3}$ quarks. Of particular importance is the actual production of h^0 in B decays. This is a one-loop process, dominated by the couplings of the Higgs bosons to the (virtual) t quark; and since all $Q = \frac{2}{3}$ quarks must couple to the same Φ_{I_0}, and be enhanced, this coupling is determined by the $J/\psi \to \gamma + H^0$ branching ratio. The branching ratio for $B \to h^0 +$ anything in the minimal model (MM) has been computed[11] from the diagrams of Fig. 1:

$$R = \left. \frac{\Gamma(B \to h^0 + \text{any})}{\Gamma(B \to e\nu + \text{any})} \right|_{\text{MM}} = \frac{|K_{tb} K_{ts}^*|^2}{|K_{cb}|^2} \frac{27\sqrt{2}}{64\pi^2} G_F m_b^2 \left(\frac{m_t}{m_b} \right)^4 \varphi \equiv R_0. \tag{12}$$

φ is a phase-space factor, normalized to 1 in the limit $m_H^2, m_c^2 \ll m_b^2$. The K_{AB} are elements of the Kobayashi-Maskawa (KM) matrix; the ratio in (12) is very close to 1.[12] Then

$$R_0 \simeq 1.5 \times 10^{-5} (m_t/m_b)^4. \tag{13}$$

To work out the modification of (12) in the many-Higgs-doublet schemes, one has to work out the couplings of the charged Higgs bosons to the $Q = \frac{2}{3}$ and $Q = -\frac{1}{3}$ quarks. For the unphysical "gauge" Higgs bosons, the couplings are the same as in the minimal standard model. For the physical $H^\pm$, the dominant (enhanced) term is

$$\mathcal{L}_{q,H^\pm} \simeq \frac{g}{2\sqrt{2}} \frac{1}{m_W} \frac{1}{\beta} \sum_I{}' \Omega(H^\pm)_{I_0 I} H_I \bar{u}^A \times m_u^A (1 - \gamma_5) K^{AB} d^B + \text{H.c.} \tag{14}$$

The second, eighth, and tenth diagrams in Fig. 1 include three Higgs-fermion couplings, and have the possibility of $1/\beta^3$ enhancement. After taking into account $1 \pm \gamma_5$ factors and the approximation $m_s \ll m_b$, only a (convergent) part[11] of the second

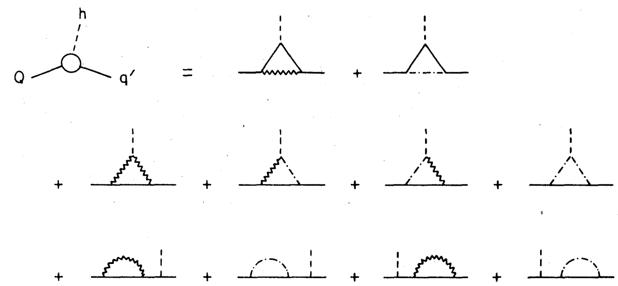

FIG. 1. Feynman diagrams for $Q \to h + q'$. Solid line, quark; wiggly line, vector boson $W^\pm$; dashed line, physical neutral Higgs scalar; dash-dotted line, charged Higgs boson.

diagram contributes a $1/\beta^3$ enhancement. Then, in the limit $\beta \ll 1$, the enhancement factor of (12) is computed to be

$$R = \left(\frac{1}{\beta^3} \right)^2 \sum_I{}' [\Omega(H^\pm)_{I_0 I} f(m_I^2)]^2 R_0, \tag{15}$$

where the m_I are the unknown masses of the physical charged Higgs bosons. Numerically, $f(m_I^2)$ ranges from 0.3 for $m_I^2 \ll m_t^2$ to 0.1 for $m_I^2 = m_t^2$. Thus, with the assumption that the masses of the $H^\pm$ are not greater than m_t [plausible if $\xi(2.2)$ is a Higgs scalar with $m_h \sim \beta m_H$], one has a lower bound on R:

$$R \gtrsim [\beta^{-3} f(m_t^2)]^2 \sum_I{}' [\Omega(H^\pm)_{I_0 I}]^2 R_0,$$
$$\sum_I{}' [\Omega(H^\pm)_{I_0 I}]^2 = \sum_I [\Omega(H^\pm)_{I_0 I}]^2 - \beta^2 = 1 - \beta^2 \simeq 1. \tag{16}$$

The primed sum explicitly excludes the contribution of the unphysical charged Higgs (whose contribution is included in R_0). Its coupling is known; it is the same as in the minimal model. Because of the b^{-3} dependence of (16) [see (7)], one can contemplate an experimental search for the semiexclusive decay mode

$$B \to h^0 + \text{anything}; \quad h^0 \to K^+ K^-. \tag{17}$$

The combined branching ratio for (17) is

$$\bar{B}B \equiv B(B \to h^0 + \text{any}) B(h^0 \to K^+ K^-)$$
$$= R \times B_{e\nu} \times b \gtrsim (r^3/b^2) f^2 R_0 B_{e\nu}. \tag{18}$$

Note that if $b \equiv B(h^0 \to K^+ K^-)$ is decreased, this combined branching ratio increases. To estimate a bound on b, Haber and Kane[13] have considered, in addition to $s\bar{s} \to K^+ K^-, K^0 \bar{K}^0$, also $s\bar{s} \to K^* \bar{K}^*$ with a factor of 3 for spin counting for each of the two charge combinations and a factor of $\frac{2}{3}$ for phase space, arriving at the conservative bound $b \lesssim \frac{1}{6}$. The existence of multiparticle final states, $K\bar{K}+\pi$'s, further decreases this.

Then, for $r>1.3$ (r_c), $1/b>6$, $f>0.1$ ($m_I<m_t$), $R_0>4\times10^{-3}$ ($m_b=4.9$ GeV, $m_t>20$ GeV), and[14] $B_{e\,\nu}=0.12$, we have $\bar{B}B>5\times10^{-4}$, which is below the sensitivity of the CLEO search. However, this lower bound is the product of four lower bounds (and very sensitive to m_t) and hence is quite an improbable value. If we take $m_I\approx m_t>30$ GeV, we obtain $\bar{B}B>2.4\times10^{-3}$, and if we also use $r_n=2.8$, rather than r_c, we obtain $\bar{B}B>24\times10^{-3}$. Thus it seems unlikely that (18) would be as small as 3×10^{-3}, the 90%-C.L. upper bound of the CLEO search,[15] if $\xi(2.2)$ were h^0.

Finally, I have examined the contribution of the enhanced Higgs couplings to the $Q=\frac{2}{3}$ quarks to the neutral flavor-changing processes $K_L\rightarrow\mu\bar{\mu}$ and m_L-m_S. I find the bounds to be less stringent than (18). Although the calculations are somewhat involved, there are two essential ingredients. The calculation of the "short distance" (quark level) contributions to $K_L\rightarrow\mu\bar{\mu}$ and m_L-m_S gives functions of x, where $x=m^2/m_W{}^2$, with $m=m_c$ or m_t. After the Glashow-Iliopoulos-Maiani mechanism is effective, these functions are $O(x)$ for small x. The sum of diagrams which contribute the leading enhancement are $O(x^2)$. This kills the contribution of enhanced $H^\pm$ coupled to virtual c quarks. For the t-quark contribution, the extra power of $m_t{}^2/m_W{}^2$ may not be sufficient to kill the enhancement, but the relevant product of KM factors has recently been determined[12] to be substantially smaller than previously known ($|K_{td}K_{ts}{}^*|<2\times10^{-3}$).

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